

Inflation from string field theory

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Abstract. In the framework of string field theory (SFT) a setting is considered, where the closed string dilaton is coupled to the open string tachyon at the final stage of an unstable brane or brane-anti-brane pair decay. The structure of non-locality leads to interesting inflationary scenarios. Concretely, we show that this configuration can lead to viable inflation by means of the dilaton becoming a non-local (infinite-derivative) inflaton. We obtain (i) a class of single field inflation with universal attractor predictions at $n_s \sim 0.967$ with any value of $r < 0.1$, where the tensor to scalar ratio r can be solely regulated by parameters of the SFT; (ii) a new class of two field conformally invariant models with a peculiar quadratic cross-product of scalar fields. In particular, we analyze a specific case where a spontaneously broken conformal invariance leads to Starobinsky like inflation, plus creating an uplifted potential minimum which accounts for vacuum energy after inflation.

Keywords: Cosmology of Theories beyond the SM, Inflation, string theory and cosmology

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1 Introduction

Primordial inflation is a compelling paradigm describing the early Universe. This is manifest through convincing observational data [1]. The end of inflation is characterized by primordial perturbations which are eventually responsible for the structure formation in the Universe. Their characteristics, namely, spectral tilts and the ratio of tensor to scalar power spectra r , have been recently measured. r has a well-established upper limit¹ $r < 0.1$ at 95% confidence level from Planck 2015 [1, 3], whereas the scalar tilt is most precisely measured as $n_s = 0.968 \pm 0.06$ at 95% confidence level. The CMB power spectra are so far found to be very much adiabatic, scale invariant and Gaussian [1, 4], supporting thereby $f(R)$ or single field inflation models. However, the physical nature of the inflaton is not certain [5, 6]. Among a broad class of models, the Starobinsky model based on the $R + R^2$ gravity modification and the Higgs inflation [7–9] occupy a privileged position, with practically equal predictions in the (n_s, r) plane

$$n_s = 1 - \frac{2}{N} \quad , \quad r = \frac{12}{N^2} \quad , \quad (1.1)$$

where N is the number of e -foldings before the end of inflation. For the expected value $N \approx 60$ the above predictions match very well the current observational values and constraints.

During the last years there have been many attempts to embed the inflationary picture into low energy effective theories derived from a more fundamental approaches, such as string theory and supergravity (SUGRA) [10–14]. Furthermore, the predictions in (1.1) provided a special stimulus to inflationary model building. More precisely, flat potentials of the following form

$$V \sim \left(1 - e^{-\sqrt{2/3}B\varphi}\right)^{2n} \quad . \quad (1.2)$$

became successful candidates for the description of inflation and were realized in various theories [5, 6, 10]. The parameter B in the above potential can lead to the universal attractor predictions with any value of $r < 0.1$

$$n_s = 1 - \frac{2}{N} \quad , \quad r = \frac{12B}{N^2} \quad . \quad (1.3)$$

¹ $r < 0.07$ considering the recent results of BICEP2/Keck Array [2].

The parameter B regulates the value of r and has different origins in different inflationary models. In particular, such a parameter arises in $\mathcal{N} = 1$ SUGRA models, as well in no-scale [15, 16], α -attractor [17, 18], higher dimensional Starobinsky [19], $f(R)$ formulation in minimal SUGRA [20, 21], $f(R)$ inflation realized with an auxiliary vector field [22] models. Inflation in SUGRA has an extra advantage, addressing the scale of supersymmetry (SUSY) breaking and a possible relation with the LHC physics [16, 23–25]. Moreover, the α -attractor models are also successful in assisting the discussion regarding the dark energy (DE) by uplifting the inflaton potential at the minimum after inflation [25–27]. A generic structure of Kähler potentials in SUGRA suitable for inflation and a possible connection to the open/closed string theory were studied in [28]. In addition, models inspired from string theory were successful in confronting Planck data [11, 12, 14]. Namely, models of axion monodromy [29–31], DBI inflation [32], p -form inflation [33]. Recently, universal attractor inflation from fibrous strings has been established in [34]. We notice also that even though inflationary model building helps shaping the more fundamental theories [12, 35], it is not enough to consider just the (n_s, r) parameters in order to single out a proper model. For example, Starobinsky and Higgs inflation models seem to differ only at the reheating stage [36]. A recent Bayesian study of models of single field slow-roll scenarios was performed with a special account to the current constraints on reheating from Planck 2015 and BICEP2/Keck Array [5, 37].

According to the present observational data, the Hubble parameter during inflation can be as large as 10^{15} GeV, suggesting the scale of inflation to be of the order of $M_I > 10^{15}$ GeV. These energy scales are acceptable in string theory, argued to play a crucial role [35]. Accounting string theory as a key player in cosmological inflation may mean employing string field theory (SFT) [38], which is necessary in building the potentials of string excitations. This is essential because to implement inflation an effective a scalar degree of freedom (the inflaton) is needed, with an appropriate potential, such that all the required characteristics (slow-roll parameters) are eventually met. It is the purpose of the present paper to demonstrate that natural models of inflation arise in the framework of SFT. We therefore name them as models of “SFT inflation” for brevity.

The inflationary process within SFT arises from the presence of the dilaton of the closed string and the tachyon of the open string. The dilaton is an integral part of the zero-mass spectrum of closed strings.² In our SFT scenario we consider open and closed strings low mass-level Lagrangian coupled through metric and dilaton. We observe that the open string tachyon upon coupling to the massless sector of closed strings generates an interesting framework for an inflationary model. The unstable brane (or brane-anti-brane pair) present in the system is known to decay, which is associated with the open string tachyon condensation (TC) [39]. Being more precise, the Sen conjecture about TC i.e., the compensation of the brane tension by the negative vacuum energy of the tachyon in the minimum of its potential, was proven using Minkowski as the embedding space-time for strings. It was established that the TC process does not require a dynamical departure from Minkowski background. This is supported by explicit papers [40, 41] and related studies which used the pure SFT framework. We should note that the rolling tachyon process does not have to be necessarily associated with an inflation or other space-time dynamics. Furthermore, models which treat the open string tachyon as the inflaton are often effective phenomenological constructions without a computational support in the SFT framework.

²However, considering only the closed string effects would not yield satisfactory SFT inflation as we will point out in the technical part of the paper.

Moreover, the present day understanding of inflation from the point of view of collected CMB data significantly favours models where the inflaton is coupled non-minimally to the Ricci scalar in the action. For example, neither DBI-type models [42, 43] nor p -adic inflation model [44] have this feature. The dilaton on contrary is naturally coupled to the Ricci scalar and thus an obvious candidate for the inflaton. In order to achieve inflation via dilaton, we assume that the string scales and energies of the brane tension are higher than the scale of inflation. We thus greatly expect, given that a brane decay should happen, inflation to start after it. It will be shown in the course of this paper that not all of the brane tension may be compensated as long as we consider the space-time dynamics implying that the space-time is curved with higher order closed string couplings near the end of the TC. As a consequence, the process of TC in open strings being essentially a non-perturbative and coupled to closed strings, contains some properties of strings interaction including infinite derivative non-locality. We stress that non-locality in the form of analytic functions of the covariant d'Alembertian is a characteristic feature of SFT and exactly the SFT non-locality will prove crucial to the concrete models of inflation we will construct.

We show that the infinite derivative non-locality brings the following two important scenarios to the inflation process. On the one hand it can produce a single field scenario in which an analog of the parameter B arises and regulates r . This parameter depends on the original structure in SFT. On the other hand a two field scenario with a conformal symmetry can also arise from SFT, which, as we show, is equivalent to the Starobinsky inflation accompanied with a non-trivial uplifting of the inflaton potential at the minimum. Notice that early attempts of considering inflation in models with infinitely many derivatives were performed in [44, 45].

In Section 2 we explain the basic concepts of SFT which are relevant for the present paper. In Section 3 we obtain an explicit inflationary framework which encapsulates the ideas of Section 2. In Section 4 we then subsequently present an effective model of inflation based on one scalar field. This model leads to a regulated value of r . In Section 5 we afterwards explore an effective two-field model of inflation. This model can be conformally invariant and leads to a positive vacuum energy value after inflation. In Section 6 we summarize and discuss open questions which follow from our SFT inflation setting.

2 Brief overview of SFT and tachyon condensation

In generic words SFT is an off-shell description of interacting strings [38, 46–50]. It describes a string by means of a string field Ψ . This object is a shorthand for encoding all the string excitations in one instance. The corresponding action for open string field³ can be written as

$$S = \frac{1}{g_o^2} \left(\frac{1}{2} \int \Psi \star Q\Psi + \frac{1}{3} \int \Psi \star \Psi \star \Psi \right) \quad (2.1)$$

where \star and \int are Witten product and integral for string fields respectively. Q is the BRST charge. The first term clearly corresponds to the motion of free strings while the second term represents the interaction. The second term is the three-string vertex responsible for the non-perturbative physics. g_o is the open string coupling constant, it is dimensionless.

It has been understood [39, 53–56] that the tachyon of open strings is responsible for the decay of unstable D -branes or D -brane-anti- D -brane pairs. The corresponding process

³An action for a closed SFT can be written only in a non-polynomial form, even for the bosonic strings [51, 52].

is the condensation of the tachyon (TC) to a non-perturbative minimum. Upon the TC the unstable brane (or pair) decays. It is the cornerstone of Sen's conjecture regarding TC that the depth of the tachyon potential minimum is exactly the tension of an unstable brane to which the string is attached to. The decay of a brane represents a configuration in which open strings must not exist, because the brane, to which they were attached, has decayed [57, 58]. This being said, we follow Sen's conjecture, which prescribes the disappearance of open string excitations. The latter phenomenon of open strings extinction can be formalized as follows in the field-theoretical language. Given a field φ the following quadratic Lagrangians are non-dynamical

$$L = -m^2\varphi^2 \text{ or } L = \varphi e^{\gamma(\square)}\varphi \quad (2.2)$$

The left Lagrangian is clearly a mass term without any dynamics. In the right Lagrangian, \square is the space-time d'Alembertian and γ is an entire function. Although it may look like \square produces dynamics as it is a differential operator, as long as we require that the function in the exponent is an entire function, the whole exponent has no eigenvalues as an operator. This means that the inverse of such an exponent gives no poles in the propagator and effectively we have no dynamics at all.

We further notice that the right Lagrangian in (2.2) is an essentially non-local Lagrangian. It is obviously non-dynamical on the quadratic level and as long as the field φ is alone. However, novel and unusual effects can be generated upon coupling to other fields or in the non-linear physics [40, 41, 44, 59, 60].

The essence of SFT is that as long as a string interaction is involved then the non-locality of the above type emerges. Technically, we can understand this as follows. Strings are extended objects by construction. When a field-theoretic model describes strings, this property of an extended object is encoded in the non-locality of interactions. SFT straightforwardly creates vertex terms of the form

$$\sim \left(e^{\alpha'\square}\varphi_1\right)\left(e^{\alpha'\square}\varphi_2\right)\left(e^{\alpha'\square}\varphi_3\right) \quad (2.3)$$

Here α' is the string length squared (which may be different from the inverse of the Planck mass squared). As it will become obvious in the course of this paper non-locality proves crucial in constructing SFT based cosmological models.

Computing any process in SFT leads to much more complicated results than presented above. TC is not an exclusion. Schematically, to describe the TC we should first compute an effective action in which all massive modes of a string with positive mass square are integrated out. Upon this computation a non-local interaction of several tachyons arise. The non-local operators are not just identical exponents but rather algebraic combinations of them. This effective action is enough to test Sen's conjecture for both, depth of the potential and absence of dynamics at the bottom of the potential. It is worth noting that actual computations in SFT are indeed difficult and technical performed by means of a level truncation scheme (i.e. including only fields up to a given mass m and the next iteration includes fields up to mass $m + 1$, etc. [61]). This scheme was proven to be convergent [61].

It is sufficient for the purposes of the present paper only to note that upon lengthy computations [38], the quadratic Lagrangian of the open string tachyon \mathcal{T} near the vacuum is non-dynamical of the form

$$L_{\mathcal{T}} = -\frac{T}{2}v(\square, \mathcal{T}). \quad (2.4)$$

For zero momenta, i.e. when $\square = 0$ the resulting $v(o, \mathcal{T})$ is exactly the tachyon potential. The dependence on \square is analytic and being linearized near the vacuum value of field \mathcal{T} it

produces

$$L_\tau = -\frac{T}{2} \frac{v''(\mathcal{T} = \mathcal{T}_0)}{2} \tau e^{\gamma(\square)} \tau, \quad (2.5)$$

with some entire function $\gamma(\square)$. The coupling T is nothing but the tension of the unstable D -brane given as

$$T = \frac{1}{2\pi^2 g_o^2 (\alpha')^{\frac{p+1}{2}}}, \quad (2.6)$$

where α' is the string length squared and p comes from the dimensionality of the Dp -brane. Thus, as expected for a 3-brane, T has a dimension length⁻⁴ and the tachyon field τ is dimensionless.

3 From SFT towards inflationary framework

In this Section we propose an action that describes the most general low mass level couplings of open and closed strings. Being precise, our action absorbs in a nut shell the description of open string tachyon coupled to the closed string dilaton at the final stages of brane decay i.e., TC. We show that such an action quite generically accommodates a positive (de Sitter) as well as a negative (Anti-de Sitter) curvature backgrounds. Given a de Sitter (dS) phase is possible we then investigate the possibility of inflation in this regime.⁴

First, we briefly describe the current context, which is particularly relevant for a open and closed SFT. We then demonstrate how a set up containing couplings between open and closed strings can lead to a cosmologically interesting scenarios, which in our case correspond to cosmic inflation.⁵ We stress that our construction not only abide to the current knowledge of SFT but in addition presents an interesting direction for the future investigations. Restricting to the physics of early Universe we show by explicit computation that viable inflationary scenarios can be naturally drawn from the structure of non-locality in our model.

We work with $(1 + 3)$ -dimensional space time assuming it is possible to organize a dimensional reduction with all moduli fields stabilized. An impact of such a compactification can be absorbed in the overall action normalization. We set the metric signature $(-, +, +, +)$, small Greek letters are the fully covariant indexes.

From closed SFT, the graviton and dilaton part of action is given by [48, 63]

$$S_c = \int d^4x \sqrt{-g} \frac{M_P^2}{2} e^{-2\phi} (R + 4\partial_\mu \phi \partial^\mu \phi). \quad (3.1)$$

Here M_P is the reduced Planck mass such that $8\pi G = \frac{1}{M_P^2}$, with G being the Newtonian constant. The dilaton field ϕ is dimensionless. Notice that it is the correct sign for the dilaton kinetic term as it appears in a closed string spectrum. Action (3.1) is the zero mass level of the closed strings. We can in principal add a p -form but it enters the action quadratically and we put it to zero using its equations of motion. Direct SFT based computation can be done to get the first orders of the terms in the latter action [64, 65].

We however do not include neither the closed string tachyon, nor any potential for the dilaton. Closed string tachyon, even though it is in the spectrum of closed strings, seems to condense to a point where the value of the field is infinite but the potential is zero (not

⁴We are mostly interested in the spin-0 (scalar) sector as it is crucial for determining inflationary parameters. The spin-2 (tensor) sector of our models turns out to be just as in GR.

⁵We note here the late time cosmological application of such a system was considered in [62].

only its derivative) [66]. Additionally it was shown in [66] that this vacuum is background independent exactly due to the fact that the field takes an infinite value. In such a way, a closed string tachyon does not contribute to our consideration of subsequent inflation. Regarding the dilaton potential, it was suggested in [67] that apparently no dilaton potential is generated. This claim finds supporting computations in the same paper and this is known as dilaton theorem.

Considering the open SFT sector we immediately make use of the result of Section 2, i.e., formula (2.4) which is relevant to describe the open string effects close to the end of an unstable D -brane decay.⁶ Moreover, if we couple it to dilaton field in a usual way supported by the linear dilaton conformal field theory on the string world-sheet (see for instance [62]), we thus obtain

$$S_o = -\frac{T}{2} \int d^4x \sqrt{-g} e^{-\phi} [v(\square, \mathcal{T}) + 1] \quad (3.2)$$

The unit term represents the brane tension. This would exactly compensate the value of the potential at the minimum in a pure open SFT and in Minkowski background where all the computation regarding the Sen's conjecture were done in a standard SFT approach. However, since the value of the tachyon field in the minimum of the potential is finite, the minimum should be background-dependent. This means that in a curved background the energy may not (and most likely will not) be compensated exactly. Moreover, we notice that the power exponent of dilaton is 1 here and not 2 as in the closed string part of the action, as it is suggested by the open SFT in a linear dilaton background [62].

Proceeding with a minimal gravitational coupling of (3.1) and (3.2) we get

$$S = S_c + S_o = \int d^4x \sqrt{-g} \left[\frac{M_P^2}{2} (\Phi^2 R + 4\partial_\mu \Phi \partial^\mu \Phi) - \frac{T}{2} \Phi [v(\square, \mathcal{T}) + 1] \right]. \quad (3.3)$$

Here we have redefined the dilaton field as $\Phi = e^{-\phi}$. The dilaton gravity on its own is a well developed subject already for a long time. See, for instance, [69] for a review. A careful but quick analysis immediately shows that this latter action is not enough. We can easily see that the Minkowski background is the only option here if there is an exact compensation of the tension of the initial D -brane by the tachyon energy at the bottom of the potential and the dilaton is a constant. Indeed, varying with respect to Φ and seeking for a constant dilaton solution (which cannot be zero as the true dilaton is $\phi = -\log(\Phi)$) we get

$$M_P^2 R = \frac{T}{2} [v(\square, \mathcal{T}) + 1].$$

This latter equation together with the trace of Einstein equations gives rise to the result that the brane tension must compensate exactly the tachyon potential value in the minimum and consequently we are left with $R = 0$. This further yields $R_{\mu\nu} = 0$. As a consequence we conclude that the set up of linear dilaton conformal field theory articulated in (3.3) is not suitable for producing an inflation in which case we essentially require a presence of dS background.

However, going further and including higher order non-minimal coupling terms of open and closed strings, we will see that constant curvature backgrounds are possible. Such terms may arise from a number of sources:

⁶To avoid confusions we notice that this is in no way the so called Vacuum SFT (VSFT) [68] but rather a linearization of the space-time action derived in *perturbative* SFT near the bottom of the tachyon potential. VSFT on contrary is a whole new construction involving a different BRST operator.

- Once the dilaton is not linear the “linear dilaton” analysis would not work. New interactions will be generated since the BRST algebra of the primary fields will get modified.
- Open-closed string interactions in general contain higher vertexes beyond the action above. These contributions generate new vertexes involving graviton, dilaton and open string tachyon.
- The so called “marginal deformation” [65] excitation in the closed strings. This operator is also of a weight zero but in fact is non-dynamical at a low-level considerations. However, its exclusion by equations of motion will generate additional terms to an effective action as well.

Taking into account the above points we therefore generalize the effective action (3.3) to the following form

$$S = \int d^4x \sqrt{-g} \left[\frac{M_P^2}{2} (\Phi^2 R + 4\partial_\mu \Phi \partial^\mu \Phi) - \frac{T}{2} \Phi \sum_{n=0}^{\infty} \Phi^n v_n(\square, \mathcal{T}) \right]. \quad (3.4)$$

where $v_0 = v(\square, \mathcal{T}) + 1$. Even though the whole next Subsection is devoted to explaining this action, it is worth noting here that even though we explicitly introduce the dilaton potential terms, this does not contradict the above mentioned dilaton theorem [67]. This is because the absence of the dilaton potential was argued and computationally verified only in a closed string sector without a coupling to open strings. Moreover, this was done only for the lowest mass levels of string excitations and when the underlying conformal field theory is built in a Minkowski bulk space-time. We clearly go beyond of all these limitations. In particular we note that paper [70] explicitly argues that violations of the dilaton theorem must occur once open-closed strings coupling is considered.

3.1 Retrieving a de Sitter setup from action (3.4)

In this Subsection we explain in a series of steps how action (3.4), which was formulated generalizing the coupling of open and closed strings cf. (3.2), allows to retrieve a satisfactory framework for inflation. This will emerge from (3.12) and (3.13), subsequently elaborated in detail in Sections 4 and 5. To have inflation we should have a reasonably approximated dS phase of evolution. It should indeed be not exactly dS, to have an exit from inflation. The dilaton itself without a self-coupling does not provide a dS solution as long as we stick to a stabilized (constant) dilaton solution.

Our first step is therefore to show that action (3.4) indeed may have a constant curvature (in particular dS) background solution with a constant dilaton field. Varying (3.4) with respect to the metric, dilaton and open string tachyon we can show that the following configuration is a solution

$$\Phi = 1, \quad \mathcal{T} = \mathcal{T}_0, \quad g_{\mu\nu} \text{ is dS with } R = R_0 = 2 \frac{T}{M_P^2} \sum_n v_{n,0}, \quad (3.5)$$

together with the following relations fulfilled

$$\sum_n v'_{n,0} = \sum_n v_{n,0} (3 - n) = 0, \quad (3.6)$$

where prime is the derivative with respect to an argument and the subscript 0 means that the function is evaluated at $\mathcal{T} = \mathcal{T}_0$. Φ is the dilaton and its particular value is irrelevant as long

as it is finite. We will pay the special account to the question how generic such configurations arise in SFT in a separate forthcoming study [71]. Notice that we need these latter relations to be satisfied only approximately and not exactly (even though it is possible) as we need a nearly dS for the exit from inflation. We recall in passing that just a single component v_n ends up with necessity with a Minkowski space-time. Thus, in order to generate dS space-time we need at least two terms with different powers of Φ in the action. Moreover we notice that the generality of our construction implies that an appearance of AdS space-time in which the quantization of strings is well-defined [72] also requires dilaton potential terms like in the SFT motivated action (3.4).

We stress that our main goal is the retrieval of satisfactory inflation and subsequently computation of inflationary observables. Our second step therefore requires validating the quadratic variation of our background action (3.4). This in turn allows us to consider only the so called equivalence class of actions which yield the same physics upon quadratic variation.⁷ To make this statement clear let us therefore consider the quadratic variation of action (3.4) in detail. Notice that we are mostly interested in scalar perturbations as they are directly affected by inclusion of new scalar fields in the consideration.

We thus proceed by writing down a quadratic variation of action (3.4) for scalar modes around the background (3.5). These modes are $\varphi = \delta\Phi$, trace of the metric perturbations h (we define $\delta g_{\mu\nu} = h_{\mu\nu}$, $h = h^\mu_\mu$) and $\tau = \delta\mathcal{T}$. A quadratic variation of action (3.4) can be written as two parts in the following way

$$\delta^{(2)}S = \delta^{(2)}S_{M_P^2} + \delta^{(2)}S_{int} \quad (3.7)$$

Generically different spins do not mix in the quadratic action i.e., tensor modes do not mix with scalar modes. Therefore the first part of the quadratic action reads

$$\delta^{(2)}S_{M_P^2} = \int d^4x \sqrt{-g} \frac{M_P^2}{2} \left[\varphi^2 R_0 + 4\partial\varphi^2 - \frac{3}{32}h \left(\square + \frac{R_0}{3} \right) h - \frac{3}{2}\varphi \left(\square + \frac{R_0}{3} \right) h \right], \quad (3.8)$$

describes the propagation of scalar (i.e., dilaton) and tensor degrees of freedom. From the above action we can exclude h by its equation of motion. Thanks to the fact that differential operators acting on h and φ are identical, we thus have $h = -8\varphi + h_{hom}$ where $(\square + R_0/3)h_{hom} = 0$. Substituting this h back in the quadratic action yields

$$\delta^{(2)}S_{M_P^2} = \int d^4x \sqrt{-g} \frac{M_P^2}{2} \varphi (2\square + 3R_0) \varphi. \quad (3.9)$$

The second part of the quadratic action after a Taylor expansion of the tachyon potential $v(\square, \mathcal{T})$ around $\mathcal{T} = \mathcal{T}_0$ reads

$$\delta^{(2)}S_{int} = -\frac{T}{2} \int d^4x \sqrt{-g} \sum_n \left[(n+1)n\varphi^2 v_{n,0} + nv'_{n,0}\varphi f(\square)\tau + \frac{v''_{n,0}}{2}\tau e^{\gamma(\square)}\tau \right], \quad (3.10)$$

which agrees with (2.5) accounting the fact that the open string tachyon on its own is not dynamical. In view of this, the function γ in the exponent must be an entire function but

⁷This will be made clearer in Sections 4 and 5.

the operator $f(\square)$ may have eigenvalues. Excluding τ by its equation of motion is dictated by $\tau = -\frac{\sum_n (nv'_{n,0})}{\sum_n v''_{n,0}} f(\square) e^{-\gamma(\square)} \varphi$. Substituting this back into action (3.10) yields

$$\delta^{(2)} S_{int} = -\frac{T}{2} \int d^4x \sqrt{-g} \varphi \left[\sum_n ((n+1)nv_{n,0}) - \frac{(\sum_n nv'_{n,0})^2}{2\sum_n v''_{n,0}} f(\square)^2 e^{-\gamma(\square)} \right] \varphi. \quad (3.11)$$

Metric fluctuations are irrelevant in this term as long as the background values of scalar fields are constant.

It is clear from the above formulae that higher curvature corrections are not relevant for us. Indeed, suppose there is a term in the action like

$$\sqrt{-g} \Phi^2 R^2$$

Such a term would produce contributions to h^2 and φh but as long as our background has constant scalar curvature and constant dilaton field the final effect of such an additional term would be just renormalization of constants in action (3.9). We see that in effect both, the spin-0 excitation of the metric and the dilaton field are combined into one joint scalar mode. Again, we can show by explicit computation that including other interactions, like for instance

$$\sqrt{-g} \Phi^2 R^2 w(\square, \tau),$$

will result in the same net result when all but one scalar fields can be excluded by equations of motion which finally results in a single scalar (non-local) excitation.

Our third step follows and it will be a bit long. We establish here why our proposal (3.4) provides a framework to generate a dS background and describe inflationary effects which require the second variation of the action around such a background. We recall here that the open string sector contains only the tachyon, since higher mass fields have been integrated out, in the course of the brane decay consideration. Thus in the nearly dS phase when the scalar curvature does not change considerably, we get from (3.7), (3.9) and (3.11) the following action that describes the propagation of scalar perturbations

$$\delta^{(2)} S_{M_P^2} = \frac{1}{2} \int d^4x \sqrt{-g} \varphi \mathcal{F}(\square) \varphi, \quad (3.12)$$

where

$$\mathcal{F}(\square) = M_P^2 (2\square + 3R_0) - T \left[\sum_n ((n+1)nv_{n,0}) - \frac{(\sum_n nv'_{n,0})^2}{2\sum_n v''_{n,0}} f(\square)^2 e^{-\gamma(\square)} \right]. \quad (3.13)$$

To generate inflation we must accompany our model with an appropriate potential. The linearization⁸ of (3.4) and corresponding analysis do not shed light on the form of the potential though. Rigorously speaking, a potential would follow from SFT provided we have computational abilities to extract one. At present, the state of the art of the knowledge in SFT

⁸We here note that additional contributions to scalar and tensor modes can be generated by means of adding the curvature squared corrections, like $R_{\mu\nu}^2$ or C^2 where C is the Weyl tensor. Moreover, following the recent studies performed in [73, 74] one has to pay special attention in order to maintain unitarity upon inclusion of terms which modify the Lagrangian for tensor modes beyond the Einstein's gravity. A standard minimal structure like C^2 in the action will generate a massive spin-2 ghost (see [75] for the first comprehensive study of this question). We therefore leave the full consideration as an open question.

lacks established methods to do so. In the course of this paper we will continue by assuming potentials which do not violate general principles of SFT construction. This strategy can be reversed and be used to constrain perhaps certain parameters in SFT, given we will reach eventually the ability to do such computations directly in the SFT framework.

Considering action (3.12) for a general operator function $\mathcal{F}(\square)$ we cannot convey inflationary physics straightforwardly. In general, $\mathcal{F}(\square)$ being considered as an algebraic function may have many roots. That is, equation

$$\mathcal{F}(z) = 0 \tag{3.14}$$

can have more than one solution. We name it a characteristic equation. Because of that, the propagator for the field φ will have more than one pole. As such, it is equivalent to multiple degrees of freedom. Let us therefore write a local realization of (3.12). Originally, this was done in [60] and then formalized in [76–78]. We use the Weierstrass factorization [60] which prescribes that any analytic function (we recall that SFT ensures that operators $\mathcal{F}(\square)$ are analytic functions) can be considered as

$$\mathcal{F}(z) = e^{\gamma(z)} \prod_j (z - z_j)^{m_j}, \tag{3.15}$$

where z_j are roots of the characteristic equation and m_j are their respective multiplicities. We assume hereafter that all $m_j = 1$ for simplicity. $\gamma(z)$ is an entire function and as such its exponent has no roots on the whole complex plane. It was shown in [60] that for a quadratic Lagrangian of the type (3.12) the following local equivalent Lagrangian can be constructed as

$$S_{local} = \frac{1}{2} \int d^4x \sqrt{-g} \sum_j \mathcal{F}'(z_j) \varphi_j (\square - z_j) \varphi_j \tag{3.16}$$

where prime means derivative with respect to the argument z with the further evaluation at the point z_j . It is said to be equivalent, thanks to the fact that solution for φ which can be obtained from equations of motion following from (3.12) is connected to solutions for φ_j simply

$$\varphi = \sum \varphi_j. \tag{3.17}$$

Roots z_j become the most crucial objects in classifying our model. Several comments are in order here:

- Note that roots z_j can be complex in general. One real root z_1 is the simplest situation. In this case, we have just a Lagrangian for a massive scalar. It is acceptable if $\mathcal{F}'(z_1) > 0$ in order to evade a ghost in the spectrum. The corresponding inflation scenario is considered in Section 4.
- More than one real root apparently seems not to be a promising scenario. Since the function $\mathcal{F}(z)$ is analytic (and therefore continuous), neighbouring real roots will be accompanied with $\mathcal{F}'(z_j)$ of opposite signs. In other words, one root is normal and the next to it is a ghost. In principle two real roots can render a well-defined model when the coupling to gravity is accounted. We will return to this in Section 5.
- Another option is complex roots. Actually, these will prove to be crucial as we will show in Section 5. Since $\mathcal{F}(z)$ is a polynomial, each complex root has its complex conjugate.

Considering such a pair of complex conjugate roots, we have

$$S_{pair} = \frac{1}{2} \int d^4x \sqrt{-g} [\mathcal{F}'(z_1) \varphi_1 (\square - z_1) \varphi_1 + \mathcal{F}'(\bar{z}_1) \bar{\varphi}_1 (\square - \bar{z}_1) \bar{\varphi}_1] , \quad (3.18)$$

where a bar over represents the complex conjugates. To maintain the connection with the original action (3.12) we should consider complex conjugate solutions to equations of motion, such that $\varphi = \varphi_1 + \bar{\varphi}_1$ is real. The important feature is that the quadratic form of fields is already diagonal. Introducing $\varphi_1 = \chi + i\sigma$, $z_1 = \alpha + i\beta$, $\mathcal{F}'(z_1) = c + is$ we can rewrite action (3.18) in terms of real components as

$$S_{pair} = \int d^4x \sqrt{-g} [\chi(c\square - c\alpha + s\beta)\chi - \sigma(c\square - c\alpha + s\beta)\sigma - 2\chi(s\square - s\alpha - c\beta)\sigma] \quad (3.19)$$

The above action is inevitably non-diagonal and features a cross-product of real fields $\sim \chi\sigma$. This cross-product plays a crucial role which will be discussed in Section 5. In the formulation above, one field looks like a ghost and such complex roots were considered as forbidden in early studies [79]. In this paper, however, we do have coupling to gravity. This coupling can save the situation in certain circumstances as will be explained shortly.

- More than one pair of complex conjugate roots or a general multi-field configuration seems non-trivial to include as will be discussed again in Section 5.

Below we consider two case studies and convey how the elaboration throughout this Section constituted indeed a framework to establish SFT inflation.

4 Attractor inflation from SFT

In this Section, in particular we study an effective action which is perturbatively equivalent up to quadratic order to (3.12) around dS background in the case of a single real root of $\mathcal{F}(\square)$. We then study the inflationary scenarios with a suitable choice of potentials and identify the role of SFT parameters in the corresponding inflationary predictions. We start with the following action

$$S_1 = \int d^4x \sqrt{-g} \left[\frac{M_P^2}{2} \tilde{\Phi}^2 R - \frac{A}{2} \partial \tilde{\Phi}^2 - V(\tilde{\Phi}) \right] . \quad (4.1)$$

From the equation of motion for the field $\tilde{\Phi}$ and Einstein equations a dS solution can be obtained in the form

$$\tilde{\Phi} = \tilde{\Phi}_0 \text{ and } R_0 = 4 \frac{V(\tilde{\Phi}_0)}{M_P^2 \tilde{\Phi}_0^2} = \frac{V'(\tilde{\Phi}_0)}{M_P^2 \tilde{\Phi}_0} > 0 .$$

Following the similar derivation in Section 3.1 we can write the quadratic Lagrangian for the spin-0 part $\tilde{\varphi} = \delta \tilde{\Phi}$ (i.e. scalar, which combines by virtue of equation of motion for the scalar field variation and the spin-0 metric perturbation) as

$$S_{1,0} = \frac{1}{2} \int d^4x \sqrt{-g} \tilde{\varphi} \left(6M_P^2 \left(\square + \frac{R_0}{2} \right) + A\square - V''(\tilde{\Phi}_0) \right) \tilde{\varphi} . \quad (4.2)$$

Assuming a SFT case where we arrive at (3.16) with just one component i.e., $j = 1$ (in which case of course all factors and fields are real), we can juxtapose (3.16) and (4.2) and observe the following identification

$$\begin{aligned}\mathcal{F}'(z_1) &\rightarrow 6M_P^2 + A \\ \mathcal{F}'(z_1)z_1 &\rightarrow 3M_P^2 R_0 - V''(\tilde{\Phi}_0) .\end{aligned}\tag{4.3}$$

In an inflating Universe, the gravity part $M_P^2 R_0$ dominates over V'' . As such, we see that SFT produces inflation and allows $\mathcal{F}'(z_1)z_1$ to be $3M_P^2 R_0$.

To simplify, we put $M_P^2 = 1$. The parameter A is related to the non-local operator $\mathcal{F}(\Box)$ and its value plays a crucial role for the subsequent characteristics of inflation. A conformal transformation $g_{\mu\nu} \rightarrow \tilde{\Phi}^{-2} g_{\mu\nu}$ of action in (4.1) gives

$$S_{1E} = \int d^4x \sqrt{-g_E} \left[\frac{1}{2} R_E - \frac{A+6}{2\tilde{\Phi}^2} (\partial\tilde{\Phi})^2 - V_E(\tilde{\Phi}) \right], \tag{4.4}$$

where the Einstein frame potential is $V_E(\tilde{\Phi}) = \frac{V_J(\tilde{\Phi})}{\tilde{\Phi}^4}$. A quartic potential in the Jordan frame $V_J(\tilde{\Phi}) = V_0 \tilde{\Phi}^4$ corresponds to a constant, after the conformal transformation such that $V_E = V_0$, in which case the field $\tilde{\Phi}$ would be massless. Consequently, the SFT setting gives a dS expansion with a constant Hubble parameter $H^2 \sim V_0$, if the kinetic energy of the massless field $\tilde{\Phi}$ is sufficiently small enough. Since the early Universe requires an exit of inflation⁹ after about $N = 50 - 60$ e -foldings of a quasi-dS expansion, we consider potentials beyond the quartic one. Therefore, a generic Jordan frame potential can be written as

$$V_J(\tilde{\Phi}) = \sum_{n=2}^{\infty} c_{n-2} \tilde{\Phi}^{2n}. \tag{4.5}$$

The field $\tilde{\Phi}$ can be canonically normalized as $\tilde{\Phi} = e^{-\sqrt{\frac{1}{A+6}}\tilde{\phi}}$. In terms of a canonically normalized field the Einstein frame potential (4.5) takes the form¹⁰

$$V_E = V_0 \sum_{n=2}^{\infty} \left(1 - e^{-(2n-2)\sqrt{\frac{1}{A+6}}\tilde{\phi}} \right), \tag{4.6}$$

where we have assumed $c_0 = -c_{n-1}$. Such a potential in (4.6) becomes exponentially flat in the limit $\tilde{\phi} \gg 1$ making it suitable for inflation and in principle, we can omit the higher order terms keeping only the leading contributions in (4.6). Therefore, for simplicity we consider here the Jordan frame potential of the form $V_J(\tilde{\Phi}) = V_0 (-\tilde{\Phi}^2 + \tilde{\Phi}^4)^2$ which, in the Einstein frame looks like

$$V_E = \tilde{V}_0 \left(1 - e^{-\sqrt{\frac{2}{3B}}\tilde{\phi}} \right)^2, \tag{4.7}$$

where $B = 1 + A/6$. The inflationary predictions corresponding to the potential in (4.7) are well known to be [15, 18, 19, 25]

$$n_s = 1 - \frac{2}{N}, \quad r = \frac{12B}{N^2}.$$

⁹Exit of inflation in string theory was extensively discussed in [35].

¹⁰See also recent studies based on the RG-corrections [80].

We assume 60 e -foldings, i.e. $N = 60$. For $B \approx 1$ (i.e. when $r \approx 0.0033$) we obtain $A = \mathcal{F}'(z_1) - 6 \approx 0$. For $B \approx 0$ (i.e., when $r \approx 0$) we obtain $A \approx -6$. Therefore, for the model to interpolate between $0 \lesssim r \lesssim 0.0033$ we require $-6 \lesssim A \lesssim 0$. Also we can obtain $r > 0.0033$ if $A > 0$. It is important to notice that $A > -6$ corresponds to a non-ghost scalar field as it is clear from the Einstein frame action (4.4).

We therefore conclude that provided the non-local operator $\mathcal{F}(\square)$ originating from the string field theory contains one real root, it gives a successful inflation with a universal attractor predictions of $n_s = 0.967$ and any the value of the tensor to scalar ratio $r < 0.1$. The value of r can be regulated to any value by varying the slope of the non-local function $\mathcal{F}'(z_1)$ at the position of the root z_1 .

5 Conformal models from SFT

In this Section, we study instead a broader case, an effective action which is perturbatively equivalent up to quadratic order to (3.19) around dS background. We start with the following non-standard action

$$S_2 = \int d^4x \sqrt{-g} \left[\frac{M_P^2}{2} [\tilde{\alpha} \tilde{\Phi}_1^2 - \tilde{\alpha} \tilde{\Phi}_2^2 - 2\tilde{\beta} \tilde{\Phi}_1 \tilde{\Phi}_2] f\left(\frac{\tilde{\Phi}_2}{\tilde{\Phi}_1}\right) R + \frac{A}{2} [\tilde{\alpha} \partial \tilde{\Phi}_1^2 - \tilde{\alpha} \partial \tilde{\Phi}_2^2 - 2\tilde{\beta} \partial_\mu \tilde{\Phi}_1 \partial^\mu \tilde{\Phi}_2] f\left(\frac{\tilde{\Phi}_2}{\tilde{\Phi}_1}\right) - V(\tilde{\Phi}_1, \tilde{\Phi}_2) \right]. \quad (5.1)$$

This particular form is motivated so that we can implement conformal symmetry and we will explain why it is of relevance in what follows. The non-standard ingredient is the cross-product of scalar fields. Using the equations of motion for the fields $\tilde{\Phi}_1$, $\tilde{\Phi}_2$ and Einstein equations we can obtain a dS solution in the form

$$\tilde{\Phi}_1 = \tilde{\Phi}_{1,0}, \quad \tilde{\Phi}_2 = \tilde{\Phi}_{2,0} \\ R_0 = 4 \frac{V_0}{M_P^2 I_0} = \frac{2\partial_{\tilde{\Phi}_1} V_0}{M_P^2 \partial_{\tilde{\Phi}_1} I_0} = \frac{2\partial_{\tilde{\Phi}_2} V_0}{M_P^2 \partial_{\tilde{\Phi}_2} I_0} > 0,$$

where we have defined $I(\tilde{\Phi}_1, \tilde{\Phi}_2) = [\tilde{\alpha} \tilde{\Phi}_1^2 - \tilde{\alpha} \tilde{\Phi}_2^2 - 2\tilde{\beta} \tilde{\Phi}_1 \tilde{\Phi}_2] f(\tilde{\Phi}_2/\tilde{\Phi}_1)$ and used short notations like $I_0 \equiv I(\tilde{\Phi}_{1,0}, \tilde{\Phi}_{2,0})$, $\partial_{\tilde{\Phi}_1} I_0 \equiv \partial I(\tilde{\Phi}_1, \tilde{\Phi}_2)/\partial \tilde{\Phi}_1$ at the values of fields at vacuum and so on for analogous terms. Following the similar derivation in Section 3.1 we can write the quadratic Lagrangian for the spin-0 part which contains 2 components $\tilde{\chi} = \delta \tilde{\Phi}_1$ and $\tilde{\sigma} = \delta \tilde{\Phi}_2$ (i.e. again the spin-0 metric perturbation is excluded by equations of motion)

$$S_{2,0} = \frac{1}{2} \int d^4x \sqrt{-g} [\tilde{\chi} \Delta_{\tilde{\chi}} \tilde{\chi} + \tilde{\sigma} \Delta_{\tilde{\sigma}} \tilde{\sigma} + \tilde{\chi} \Delta_{\tilde{\chi}\tilde{\sigma}} \tilde{\sigma}], \quad (5.2)$$

where

$$\Delta_{\tilde{\chi}} = \frac{M_P^2}{2} \left(\frac{(\partial_{\tilde{\Phi}_1} I_0)^2}{I_0} (3\square + R_0) + \frac{\partial^2 I_0}{\partial \tilde{\Phi}_1^2} R_0 \right) - A \tilde{\alpha} f_0 \square - \frac{\partial^2 V_0}{\partial \tilde{\Phi}_1^2}, \\ \Delta_{\tilde{\sigma}} = \frac{M_P^2}{2} \left(\frac{(\partial_{\tilde{\Phi}_2} I_0)^2}{I_0} (3\square + R_0) + \frac{\partial^2 I_0}{\partial \tilde{\Phi}_2^2} R_0 \right) + A \tilde{\alpha} f_0 \square - \frac{\partial^2 V_0}{\partial \tilde{\Phi}_2^2}, \\ \Delta_{\tilde{\chi}\tilde{\sigma}} = \frac{M_P^2}{2} \left(\frac{\partial_{\tilde{\Phi}_1} I_0 \partial_{\tilde{\Phi}_2} I_0}{I_0} (3\square + R_0) + \frac{\partial^2 I_0}{\partial \tilde{\Phi}_1 \partial \tilde{\Phi}_2} R_0 \right) - A \tilde{\beta} f_0 \square - \frac{\partial^2 V_0}{\partial \tilde{\Phi}_1 \partial \tilde{\Phi}_2}.$$

From a SFT framework we can come to (3.19), which is the case of two complex conjugate roots with the Lagrangian written in real fields. Hence, we can try to juxtapose (3.19) and (5.2). The motivation for doing this is to find a more fundamental origin for the effective model (5.1). This is, however much more involved than in the previous Section with a single field. Essentially, the most important is to establish $\Delta_{\tilde{\chi}} = -\Delta_{\tilde{\sigma}}$. On this way, we can neglect the second derivatives of the potential V . However, we must satisfy a number of constraints, namely, all parameters and vacuum fields values must be real and I_0 strictly positive. And we want to have $\tilde{\beta} \neq 0$, which we will explain why in the following. The greatly simplifying point is that we must require such an adjustment of coefficients of Δ -s only in a single point ($\Phi_1 = \Phi_{1,0}$, $\Phi_2 = \Phi_{2,0}$). On top of this we emphasize once again that we aim at retrieving a nearly dS phase, not an exact one. These requirements are generically satisfied altogether with the presence of a function $f\left(\frac{\tilde{\Phi}_2}{\tilde{\Phi}_1}\right)$ (apart from special situations which we discuss shortly). It is important that being a function of the ratio of fields it cannot spoil a possible conformal invariance.

Notice that our main purpose in this Section is to establish an effective setting which can emulate (3.19). We claim that we have such an effective model as long we can match quadratic actions for scalar modes around a dS background. We can establish a correspondence between (3.19) and (5.2) by means of the following:

- During inflationary expansion we can assume that the scalar fields varies slowly and the kinetic terms can be neglected. We are thus mainly interested in whether $\Delta_{\tilde{\chi}} = -\Delta_{\tilde{\sigma}}$ for the terms proportional to R_0 . To have this we should require

$$\frac{(\partial_{\tilde{\Phi}_1} I_0)^2}{I_0} + \frac{\partial^2 I_0}{\partial \tilde{\Phi}_1^2} + \frac{(\partial_{\tilde{\Phi}_2} I_0)^2}{I_0} + \frac{\partial^2 I_0}{\partial \tilde{\Phi}_2^2} \approx 0 \quad (5.3)$$

- We can check that even in the very simple case of $\tilde{\beta} = 0$, a non-constant function f is required to satisfy the above relation. A simple choice like

$$f = 1 + f_1 \tilde{\Phi}_2 / \tilde{\Phi}_1, \quad (5.4)$$

with just one free parameter f_1 is sufficient. Otherwise, for $f = \text{const}$ a condition $\tilde{\Phi}_{1,0} = \pm i \tilde{\Phi}_{2,0}$ arises from (5.3). Therefore to build such an effective model the function $f\left(\frac{\tilde{\Phi}_2}{\tilde{\Phi}_1}\right)$ is very useful and important. The cross-product of fields may arise for $\tilde{\beta} = 0$ but a quite involved non-polynomial function f is required.

- For a non-trivial $\tilde{\beta}$ the same function f as above in (5.4) is enough to arrange the condition (5.3). Moreover $\tilde{\beta} \neq 0$ generates a cross-product of fields.
- In complete analogy we can consider the coefficients of the kinetic terms. We have to require a non-constant function f . We note that having opposite coefficients in front of d'Alembertian operators for different fields essentially means that one of these fields is a ghost.

Recalling the SFT expressions (3.16) and (3.19), we see that the presence of a cross-product is a special feature related to a complex root of the function $\mathcal{F}(z)$ (which defines the non-local operator $\mathcal{F}(\square)$). This means that the parameter β found in (3.19) is essentially non-zero (notice that there is no a direct simple relation between $\tilde{\beta}$ and β). In the limiting case of

$\beta \rightarrow 0$, we should see the cross-product disappearing and this corresponds to $\tilde{\beta} \rightarrow 0$ in the effective model (5.1). Another way to recognize the effective model (5.1) without a cross-product of fields is to consider directly (3.16) with two specially tuned real roots. This means that these roots are related as $z_2 = -z_1$ and moreover $\mathcal{F}'(z_2) = -\mathcal{F}'(z_1)$.

To resolve the issue of a ghost in the spectrum requires an extra symmetry in order to gauge the ghost away. The most natural candidate is the conformal symmetry used in the building of similar models in [17, 18, 81]. The conformal invariance is restored in (5.1) if we assume $A = 6$. Our model without a cross-product resembles the conformal models studied in [82, 83]. We stress that the cross-product appeared for the first time in the cosmological models and we have here provided an imperative explanation through the SFT framework. In the following Subsection, inflation in a model with such a cross-product of scalar fields is analyzed.

5.1 Conformal inflation and Vacuum energy

We start with our effective action in (5.1) assuming $f\left(\frac{\tilde{\Phi}_2}{\tilde{\Phi}_1}\right) \approx \text{constant}$ during inflation

$$S_3 = \int d^4x \sqrt{-g} \left[\left(\tilde{\alpha} \tilde{\Phi}_1^2 - \tilde{\alpha} \tilde{\Phi}_2^2 - 2\tilde{\beta} \tilde{\Phi}_1 \tilde{\Phi}_2 \right) \frac{R}{12} + \frac{\tilde{\alpha}}{2} \partial \tilde{\Phi}_1^2 - \frac{\tilde{\alpha}}{2} \partial \tilde{\Phi}_2^2 - \tilde{\beta} \partial_\mu \tilde{\Phi}_1 \partial^\mu \tilde{\Phi}_2 - V_J \left(\tilde{\Phi}_1, \tilde{\Phi}_2 \right) \right], \quad (5.5)$$

where we have set $M_P = 1$ for simplicity and use the subscript J for the Jordan frame as before. Since the field $\tilde{\Phi}_1$ has a wrong sign kinetic term (assuming $\tilde{\alpha} > 0$), we can eliminate it by the choice of conformal gauge $\tilde{\Phi}_1 = \sqrt{6}$ which spontaneously breaks the conformal invariance. To obtain a consistent inflation within this model we consider the following potential

$$V_J \left(\tilde{\Phi}_1, \tilde{\Phi}_2 \right) = \frac{\lambda}{4} \left(\gamma_1 \tilde{\Phi}_2^2 + \gamma_2 \tilde{\Phi}_1 \tilde{\Phi}_2 + \gamma_3 \tilde{\Phi}_1^2 \right) \left(\tilde{\Phi}_2 - \tilde{\Phi}_1 \right)^2, \quad (5.6)$$

where $\gamma_1, \gamma_2, \gamma_3$ are arbitrary constant parameters. The potential (5.6) is motivated from [17], which we generalize here to our conformal model with a term containing the cross-product of fields. The importance of this generalization will be explained in what follows. Note that if $\tilde{\beta} = \gamma_2 = \gamma_3 = 0$, the model reduces to the conformal model without a cross-product of fields studied in [17].

Rescaling the fields as $\tilde{\Phi}_1 \rightarrow \frac{\tilde{\Phi}_1}{\sqrt{\tilde{\alpha}}}$ and $\tilde{\Phi}_2 \rightarrow \frac{\tilde{\Phi}_2}{\sqrt{\tilde{\alpha}}}$ in action (5.5) and using the gauge $\tilde{\Phi}_1 = \sqrt{6}$ we yield

$$S_3 = \int d^4x \sqrt{-g} \left[\frac{R}{2} \left(1 - \frac{\tilde{\Phi}_2^2}{6} - \frac{2\tilde{\beta}}{\sqrt{6}\tilde{\alpha}} \tilde{\Phi}_2 \right) - \frac{1}{2} \partial_\mu \tilde{\Phi}_2 \partial^\mu \tilde{\Phi}_2 - \frac{\lambda}{4\tilde{\alpha}^2} \left(\gamma_1 \tilde{\Phi}_2^2 + \gamma_2 \tilde{\Phi}_1 \tilde{\Phi}_2 + \gamma_3 \tilde{\Phi}_1^2 \right) \left(\tilde{\Phi}_2 - \sqrt{6} \right)^2 \right]. \quad (5.7)$$

We can rewrite the latter action as

$$S_3 = \int d^4x \sqrt{-g} \left[\frac{R}{2} \left[1 + \frac{\tilde{\beta}^2}{\tilde{\alpha}^2} - \frac{1}{6} \left(\tilde{\Phi}_2 + \frac{\tilde{\beta}}{\tilde{\alpha}} \sqrt{6} \right)^2 \right] - \frac{1}{2} \partial_\mu \tilde{\Phi}_2 \partial^\mu \tilde{\Phi}_2 - \frac{\lambda}{4\tilde{\alpha}^2} \left(\gamma_1 \tilde{\Phi}_2^2 + \gamma_2 \tilde{\Phi}_1 \tilde{\Phi}_2 + \gamma_3 \tilde{\Phi}_1^2 \right) \left(\tilde{\Phi}_2 - \sqrt{6} \right)^2 \right]. \quad (5.8)$$

Now performing the conformal transformation $g_{\mu\nu} \rightarrow \left[1 + \frac{\tilde{\beta}^2}{\tilde{\alpha}^2} - \frac{1}{6} \left(\tilde{\Phi}_2 + \frac{\tilde{\beta}}{\tilde{\alpha}} \sqrt{6} \right)^2 \right]^{-1} g_{\mu\nu}$ and shifting the field $\tilde{\Phi}_2 \rightarrow \tilde{\Phi}_2 + \frac{\tilde{\beta}}{\tilde{\alpha}} \sqrt{6}$, we arrive to the Einstein frame action

$$S_{3E} = \int d^4x \sqrt{-g_E} \left[\frac{R_E}{2} - \frac{\omega}{2 \left(\omega - \frac{\tilde{\Phi}_2^2}{6} \right)^2} \partial_\mu \tilde{\Phi}_2 \partial^\mu \tilde{\Phi}_2 - V_E \left(\tilde{\Phi}_2 \right) \right], \quad (5.9)$$

where $\omega = 1 + \frac{\tilde{\beta}^2}{\tilde{\alpha}^2}$ and

$$V_E \left(\tilde{\Phi}_2 \right) = \frac{9\lambda}{\tilde{\alpha}^2} \frac{\left[\gamma_1 \tilde{\Phi}_2^2 + \left(\gamma_2 - 2\gamma_1 \frac{\tilde{\beta}}{\tilde{\alpha}} \right) \sqrt{6} \tilde{\Phi}_2 + 6 \left(\gamma_1 \frac{\tilde{\beta}^2}{\tilde{\alpha}^2} - \gamma_2 \frac{\tilde{\beta}}{\tilde{\alpha}} + \gamma_3 \right) \right] \left(\tilde{\Phi}_2 - \sqrt{6} \frac{\tilde{\beta}}{\tilde{\alpha}} - \sqrt{6} \right)^2}{\left(6\omega - \tilde{\Phi}_2^2 \right)^2}. \quad (5.10)$$

If γ_i are chosen such that $\gamma_2 = 2\gamma_1 \frac{\tilde{\beta}}{\tilde{\alpha}}$ and $\gamma_1 \frac{\tilde{\beta}^2}{\tilde{\alpha}^2} - \gamma_2 \frac{\tilde{\beta}}{\tilde{\alpha}} + \gamma_3 \gtrsim 0$, we can obtain inflation with an uplifting of the potential at the minimum.

For example, let us consider a simple case with $\gamma_1 = 1$, $\gamma_2 = 2\frac{\tilde{\beta}}{\tilde{\alpha}}$ and $\gamma_3 = 2\frac{\tilde{\beta}^2}{\tilde{\alpha}^2}$, for which (5.10) reduces to the form

$$V_E \left(\tilde{\Phi}_2 \right) = \frac{9\lambda}{\tilde{\alpha}^2} \frac{\left[\tilde{\Phi}_2^2 + 6\frac{\tilde{\beta}^2}{\tilde{\alpha}^2} \right] \left(\tilde{\Phi}_2 - \sqrt{6} \frac{\tilde{\beta}}{\tilde{\alpha}} - \sqrt{6} \right)^2}{\left(6\omega - \tilde{\Phi}_2^2 \right)^2}, \quad (5.11)$$

The canonically normalized field $\tilde{\phi}$ is given by the relation $\tilde{\Phi}_2 = \sqrt{6\omega} \tanh \left(\frac{\tilde{\phi}}{\sqrt{6}} \right)$. The potential V_E in (5.11) in terms of $\tilde{\phi}$ reads

$$V_E \left(\tilde{\phi} \right) = \mu^2 \sinh^2 \left(\frac{\tilde{\phi}}{\sqrt{6}} \right) \left[\cosh \left(\frac{\tilde{\phi}}{\sqrt{6}} \right) - \frac{1}{1 + \frac{\tilde{\beta}}{\tilde{\alpha}}} \sqrt{1 + \frac{\tilde{\beta}^2}{\tilde{\alpha}^2}} \sinh \left(\frac{\tilde{\phi}}{\sqrt{6}} \right) \right]^2 + \frac{\mu^2 \tilde{\beta}^2}{\tilde{\alpha}^2 \left(1 + \frac{\tilde{\beta}^2}{\tilde{\alpha}^2} \right)} \cosh^2 \left(\frac{\tilde{\phi}}{\sqrt{6}} \right) \left[\cosh \left(\frac{\tilde{\phi}}{\sqrt{6}} \right) - \frac{1}{1 + \frac{\tilde{\beta}}{\tilde{\alpha}}} \sqrt{1 + \frac{\tilde{\beta}^2}{\tilde{\alpha}^2}} \sinh \left(\frac{\tilde{\phi}}{\sqrt{6}} \right) \right]^2, \quad (5.12)$$

where $\mu^2 = \frac{9\lambda(\tilde{\alpha}+\tilde{\beta})^2}{\tilde{\alpha}^2(\tilde{\alpha}^2+\tilde{\beta}^2)}$. In the limit $\frac{\tilde{\beta}}{\tilde{\alpha}} \ll 1$, the first line in (5.12) dominates during inflation while the second line is negligible. The potential (5.12) is always positive and in particular has a non-zero value at the minimum at $\tilde{\phi} \approx 0$. In general the shape of the potential is similar to the Starobinsky-like models in no-scale SUGRA [15]. In Fig. 1 we depict the shape of

the potential for various values of $\tilde{\beta}$. This corresponds to different values of vacuum energy (Λ) after inflation. We can see that the smaller the value of $\tilde{\beta}$, the greater the chance of approaching the plateau region of the Starobinsky model, and eventually the smaller will be the value of the vacuum energy.

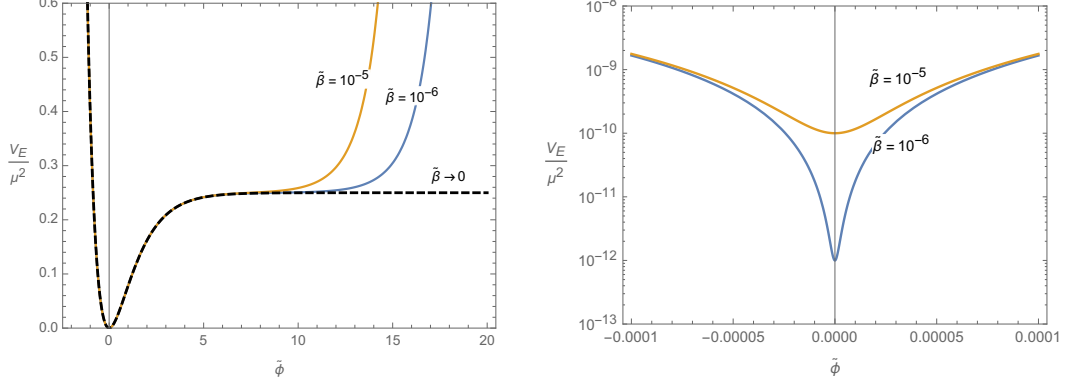


Figure 1. In the left panel we plot the potential $V_E(\tilde{\phi})$ for values of $\tilde{\beta} = 10^{-5}$, 10^{-6} and $\tilde{\alpha} = 1$. In the right panel, we depict the corresponding minimum of the potential around $\tilde{\phi} \approx 0$.

In the limit $\tilde{\beta}/\tilde{\alpha} \ll 1$, we can approximate the potential in (5.12) as (setting also $\tilde{\alpha} = 1$)

$$V_E(\tilde{\phi}) \approx \frac{\mu^2}{4} \left(1 - e^{-\sqrt{\frac{2}{3}}\tilde{\phi}}\right)^2 + \frac{\mu^2 \tilde{\beta}^2}{4} \left(1 + e^{-\sqrt{\frac{2}{3}}\tilde{\phi}}\right)^2, \quad (5.13)$$

where the first term dominates when $\tilde{\phi} \gg 1$ and leads to a Starobinsky like inflation i.e., $n_s \sim 0.967$, $r \sim 0.0033$ for $N = 60$ and the second term gives a non-zero vacuum energy at the minimum of the potential¹¹ near $\tilde{\phi} = 0$. Here $\mu \approx 2 \times 10^{-5}$ (in Planck units as we have set $M_P = 1$) which can be determined from the observed amplitude of scalar perturbations $A_s = 2.2 \times 10^{-9}$ at the horizon exit [1]. In particular $\tilde{\beta} \sim 10^{-55}$ gives a vacuum energy that reproduces the present day cosmological constant $\Lambda \sim 10^{-120}$. Therefore, we conclude that a SFT induced cross-product of the fields $\tilde{\Phi}_1$ and $\tilde{\Phi}_2$ in (5.5) naturally uplifts the inflaton potential at the minimum and possibly explain the present day dark energy (assuming it is Λ CDM).

6 Conclusions and discussion

We have shown that SFT allows various configurations of scalar fields coupled to gravity. The most important are those with a closed string dilaton and the open string tachyon. The former drives the unstable brane decay and then the non-locality of the tachyon interaction is transferred to the dilaton. The configurations get their most peculiar properties due to the non-local infinite derivative operators. We selected those formed at the linearized level as $\mathcal{F}(\Box)$. This operator in turn is a characteristic feature of SFT.

The cornerstone question is about the roots z_j of the characteristic equation $\mathcal{F}(z) = 0$. Moreover, the derivatives $\mathcal{F}'(z_j)$ play an important role. This is seen from action (3.16), which describes the evolution of scalar perturbations around a dS vacuum in SFT. Its importance

¹¹A potential of similar kind can be found in the α -attractor model where the inflaton potential was uplifted due to the effect of a SUSY breaking mechanism [25].

is obvious as inflation is a dS like expansion and all the observable quantities related to scalars can be obtained from exploring the action for linear perturbations. A very important restriction is that no ghosts must be in the spectrum. This selects two configurations of roots.

First, there is the situation with one real root z_1 accompanied with a correct sign of $\mathcal{F}'(z_1)$. In this case there is one scalar perturbative degree of freedom. Such a configuration may be given the effective model description (4.1). It is important that coefficients in front of the Einstein-Hilbert term and the kinetic term of a scalar field are independent. This translated in a non-unit B -parameter which regulates the ratio of tensor to scalar spectra r exactly as in (1.3). Therefore, we emphasize here that attractor inflation is possible to realize in the framework of SFT setup. A future more accurate detection of parameter r from CMB [84] would indicate the values of z_1 and $\mathcal{F}'(z_1)$.

Second, there was the case with two roots. They can be complex conjugate and then we should look at (3.19) which is written in manifestly real components. In this scenario, we inevitably get a cross-product of fields. Moreover, one field looks like a ghost. However, kinetic and mass terms have exactly opposite signs. This suggests that a conformal symmetry may help exorcising the ghost. Indeed, building an effective model (5.1) we have taken the conformal symmetry into account and have shown that we indeed can make use of it to remove the unwanted degrees of freedom. The cross-product of fields naturally leads to an uplifting of the potential in the reheating point. In principle one can get a similar two-field model starting with two real roots which are related as $z_1 = -z_2$ and $\mathcal{F}'(z_1) = -\mathcal{F}'(z_2)$. This latter case has no cross-product of fields and falls into the considerations of [82, 83]. The novel feature here is that the conformally invariant models with a quadratic cross-product of scalar fields appear for the first time in a cosmological setup and can be naturally explained using the non-local structure of SFT.

More generic configurations with more than two fields may have no reasonably simple effective model counterpart. This is because more than one ghost would appear. In this case quite a peculiar structure may be required in order to arrange such a configuration that it will be possible to gauge away all the ghosts. However, understanding a potential power of multifield models [83] we surely leave this as an open question. We also have skipped a case of multiple roots. It can be considered analogously but requires a more complicated formula analogous to (3.16).

At last, we point to the essential question of deriving the potentials for scalar fields directly from SFT. This is a task requiring hard computations. With our study of SFT inflation using the proposed action (3.4) with higher order couplings of open and closed strings, we will be able to see how much the potentials computed directly from SFT are compatible with inflationary requirements.

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